Gluino, wino and Higgsino-like particles without supersymmetry

Tirtha Sankar Ray, Hiroshi de Sandes, Carlos A. Savoy

Institut de Physique Théorique, CEA-Saclay,
91191 Gif-sur-Yvette, France*

Abstract

Metastable charged particles produced at the LHC can decay in the quiescent period between beam crossing in the detector leading to spectacular signals. In this paper we consider augmenting the Standard Model with gluino, Higgsino and wino-like particles without invoking supersymmetry. Proton stability is ensured by imposing a discrete Z_6 parity that fixes their possible interaction and makes them metastable in a large portion of the parameter space. We investigate the possibility of employing these fields to improve gauge coupling unification, explain dark matter, generate neutrino mass and cancel flavor anomalies. We find that the masses of these fields, controlled by the flavor anomaly relations, make them visible at the LHC.

1 Introduction

If they exist, the LHC has the power to discover possible new fermions beyond the TeV milestone, in particular, strongly interacting ones. Many publications have discussed a variety of such states, their signatures at colliders and, often, the models for physics beyond the Standard Model (SM) that suggest their being. A large amount of these studies focus on heavy quarks and leptons in fundamental multiplets of the SM. On the other hand, the supersymmetric extensions of the SM generically require gauginos, in the adjoint representations of the SM gauge group, as well as Higgsinos, electroweak doublets without lepton number. The corresponding interactions are supersymmetric transforms of the gauge ones, and of the Higgs couplings, respectively. With their signatures so fixed, these are currently the most hunted particles at the LHC except, of course, for the Higgs. It goes without saying that they are instrumental for the beautiful unification of the gauge couplings allowed in these models.

This paper discusses fermions with the same SM quantum numbers as gauginos and Higgsinos in the absence of supersymmetry. Without squarks and sleptons, they cannot mimic the supersymmetric couplings and, besides the universal gauge interactions, have only effective four fermion couplings to the standard fermions. Since they transform in real representations of the SM symmetries, their masses are expected to be near the cut-off scale, unless they are protected by some mechanism. It seems plausible to take advantage of the flavor symmetry that would allow for the quarks and leptons to have masses way below the electroweak scale to suppress the mass of these exotic states. In particular, if that symmetry is gauged, they could usefully contribute to the cancellation of anomalies generated by the SM fermions. As argued below, this scenario makes sense and allows for new fermions within the LHC reach, in spite of the relatively high limit on the flavor symmetry breaking scale imposed by data on FCNC and CP violating effects.

*Laboratoire de la Direction des Sciences de la Matière du CEA et Unité de Recherche Associé au CNRS (URA2306)

Indeed, the strong hierarchy in the measured fermion masses and mixing remains a mystery till date. An elegant resolution of this problem is achieved in the framework of Froggatt-Nielsen mechanism [1]. In the simplest form this entails the prediction of an additional $U(1)_X$ flavor symmetry group, under which the SM fermions are charged. This group, which may be gauged, is spontaneously broken at some high scale by the vev of the flavon scalar (ϕ), with a conventionally assigned flavor charge -1. At the weak scale, the coefficient of an SM operator ($\hat{O}$) is suppressed by a factor $\epsilon^{|\Delta X|}$, where ΔX is the $U(1)_X$ charge mismatch in the operator ($\hat{O}$) and $\epsilon = \langle\phi\rangle/\Lambda$, where $\langle\phi\rangle$ is the flavon vev and Λ is the cut-off. By applying this rule to quark and lepton masses, including effective neutrino mass operators, their measured patterns can be fairly reproduced by suitably choosing the flavor charges.

It is reasonable to assume that at high energies, this effective theory gives way to a renormalizable UV complete theory. In this context, the cancellation of chiral anomaly related to this $U(1)_X$ charge becomes a relevant issue. Traditionally it has been assumed that such anomaly cancellation takes place at the string scale through the Green-Schwarz mechanism. Recently it has been suggested that such anomaly cancellation can be achieved through additional exotic fermions, some of which have masses in the TeV scale [2], within the reach of the present generation colliders. To facilitate anomaly cancellation, we analyze the possibility of augmenting the SM with fermions in the adjoint representation of the SM gauge group. We find that the addition of the adjoint fields $\tilde{g}(8,1,0)$ and $\tilde{w}(1,3,0)$, and $\tilde{h}_d(1,2,-\frac{1}{2})$, $\tilde{h}_u(1,2,+\frac{1}{2})$ represent the minimal content that can, in principle, cancel all the anomalies in the theory if they are endowed with flavor charges[2]. We find that the cancellation of anomaly with this minimal content implies large $U(1)_X$ charges for these fields. And thus, in many instances, some of these fields have masses in the TeV scale.

Interestingly, this also represents the fermionic sector of the minimal SUSY extension of the SM with the exception of the hypercharge gauge fermion, the bino. A similar SM singlet field charged only under $U(1)_X$, would play no role in the cancellation of mixed anomalies and may be added to the model with obvious alterations. To make the identification more alike we further propose to expand the scalar sector to include a second Higgs doublet. We assume a Z_2^H symmetry that allows one of the Higgs to give mass to the up-type and the other to the down-type. In principle the Higgs fields can be charged under the $U(1)_X$ gauge group. When the charges assigned to the various fermions in the model are not integers, the flavon breaks $U(1)_X$ into a discrete symmetry that plays an important role for the consistency of the model. We assume it to be the so-called proton hexality [3] that forbids proton decay effective operators. It also forbids mixing between the new and SM fermions, even in presence of electroweak symmetry breaking, making the lightest new fermion metastable, or even stable, hence a dark mater candidate. In particular it forbids mixing between the Higgsino and the Lepton analogous to the R parity introduced in the MSSM. However it is consistent with lepton-wino mixing through the Higgs, providing a model for type III leptogenesis.

Of course, a crucial difference is the absence of scalars other than the Higgs fields, and of the corresponding dimension four couplings. Therefore the lightest heavy fermion decays into three SM states and becomes long-lived as the cut-off scale is high. In some instances, the phenomenology of the models somewhat mimic the split supersymmetry scenarios [4], where a striking feature is the possibility of observation of metastable gluinos [5]. In the present paper, we study the possibility of having metastable fermions in the weak scale that can be stopped in the detector, where they decay in a non-standard way at a later instant. We find that the decay pattern of the fields considered in this paper are distinct from the gluinos from the split SUSY framework and can be easily distinguished.

[2]Even if it could look inappropriate, for simplicity, we name and denote these new states just like the corresponding ones in the MSSM.

We also note that both the CMS and the ATLAS experiments at LHC [6,7] have searched for these metastable states and put initial mass limits of the fields. These limits do not severely constrain this class of models at present. However they can potentially be probed in the near future with modified search strategy.

Since this is an alternative to the supersymmetric theory discussed in [2], we refer to that paper for several details and issues that are not reproduced here. However many relations are very different because the gauginos in [2] are the real thing, and the heavy states are leptons and quarks and their scalar companions. A recent paper [8] builds non-supersymmetric versions with metastable exotic quarks and leptons.

2 The Minimal and Next to Minimal Models

Although our approach should be rather generic in the framework of gauged flavor symmetry, it will be presented here in simple cases by way of illustration. In our minimal model, we consider the SM together with the Froggatt- Nielsen model described in the introduction, augmented with adjoint fermions and two weak doublets to compensate for the anomalies related to a gauged abelian flavor symmetry. We also enlarge the scalar sector to include an extra Higgs doublet. This eases anomaly compensation. These exotic fields are quite different from the supersymmetric gauginos and Higgsinos. For example they have non trivial charges under the flavor symmetry, which among others prevent the SU(2) adjoint triplet from mixing with the SU(2) doublets, and, most importantly, they have very different couplings to quarks and leptons. As discussed below, no simple example with unification of gauge couplings was found for this minimal model. In our next to minimal model, an $I = 0$ exotic charged lepton (E) is added. This gives some additional renormalization of the hypercharge coupling improving among other things the possibility of driving gauge coupling unification around the usual GUT scale.

These ersatz particles interact with the lighter SM fields essentially through effective four fermion operators. It is well known that if higher dimensional operators are added to the SM, it predictably leads to the pitfall of rapid proton decay. This means that the cut-off of the effective higher dimensional operators are essentially pushed to the GUT scale. This situation can be evaded by considering additional symmetry in the theory that prevents operators responsible for proton decay from showing up. In the present context this can be achieved by simply considering a Z_6 symmetry, proton hexality [3], under which the SM fields are suitably charged. One can embed this discrete group into the continuous symmetry group $U(1)_X$ so that when the flavon gets a vev the flavor symmetry is spontaneously broken to the proton hexality subgroup.

For each fermion f, one can separate the integer and fractional part, $X_f = x_f + Z'_f$, where x is an integer. The fractional part of the charges remains operative below the electroweak symmetry breaking scale and prevents potentially dangerous operators in the theory. The corresponding charge assignment of the SM fields are shown in Table 1. The Higgs fields are assumed to have integer flavor charges[3] to protect the discrete symmetry below the EWSB scale. The exotic fields introduced are also charged under this discrete symmetry. The choice of these charges is almost unique as it determines the possible four fermion interaction of the exotic fields and therefore their decay patterns. We exhibit the possible choices in Table 2 consistent with Majorana masses for gauginos.

The theory is defined by the most general effective Lagrangian consistent with the SM symmetries, proton hexality and a cut-off Λ. As already noted, the various operators are modulated by the Froggatt-Nielsen factors,

[3] Actually the consideration of consistent anomaly cancellation also leads to a zero fractional flavor charge for the Higgs, as was argued in [2] and [3].

	$SU(3)_c$	$SU(2)_L$	Y	$z' \times 18$
q	3	2	$1/6$	1
u	$\bar{3}$	1	$-2/3$	-1
d	$\bar{3}$	1	$1/3$	-1
l	1	2	$-1/2$	-9
e	1	1	1	9

Table 1: Charges of SM fields, $z' = z + \frac{1}{3}Y$

	$SU(3)_c$	$SU(2)_L$	Y	$z' \times 18$
$\tilde{g}$	8	1	0	9
$\tilde{w}$	1	2	0	$9, 0$
$\tilde{h}_d$	1	2	-1	± 3
$\tilde{h}_u$	1	2	1	∓ 3
E	1	1	-1	-3

Table 2: Charges of exotic fields.

i.e., powers of ϵ equal to their flavor charges. The lowest dimension operators include the Higgs couplings to fermions that yield their masses and mixing and the four fermion interactions. The lowest dimension operators that lead to FCNC put severe bounds on the cut-off scale, as we now turn to discuss. They are also responsible for the masses of the heavy fermions and their couplings to the light ones, as will be discussed later.

By convenience, we shift the flavor charge by a fraction of the hypercharge to cancel the charge of one of the Higgses (H_u) and assign a charge x_H to H_d. Let us define the combinations of flavor charges associated to the Yukawa couplings of light quarks to the two Higgs:

$$\chi_u^{ij} = x_q^i + x_{\bar{u}}^j \,, \qquad \chi_d^{ij} = x_q^i + x_{\bar{d}}^j + x_H \,, \qquad \chi_e^i = x_l^i + x_{\bar{e}}^i + x_H . \tag{1}$$

The mass matrices of the fermions are given by, $m_f^{ij} \sim \epsilon^{|\chi_f^{ij}|} v_f$, where v_f is the vev of the corresponding Higgs field. [4] From the known quark and lepton masses and mixing, one can find the values of these Yukawa charges as a function of $tan\beta = v_u/v_d \sim \epsilon^{-x_t}$.[5] For $\epsilon \sim .2$ a reasonable mass matrix is obtained if we impose the following constraints: $\chi_u^{ii} = (8,4,0)$; $\chi_d^{ii} + x_t = (7,5,3)$ and $\chi_e^{ii} \pm x_t. = (\pm 7, \pm 5, \pm 3)$. Many other examples can be found in the literature, see, *e.g.*, [9].

The effective neutrino mass matrix can be generated after integrating out the physics above the cut-off, which would result into a dimension five operator $(LH_u)^2$ in the effective Lagrangian. If the wino has half-integer flavor charge as the leptons, it can mediate Type III seesaw [10] through its coupling $\bar{\tilde{w}} H_u L$ to the lepton doublets. Integrating out the wino with mass $\epsilon^{|\chi_{\tilde{w}}|}\Lambda$, where $\chi_{\tilde{w}} = 2x_{\tilde{w}}$, one gets a second contribution to neutrino masses. In order to ensure a neutrino mass matrix with a modest hierarchy, one assumes that the $x_l^i = x_l$ is the same for all three leptons. The order of magnitude of final neutrino masses are,

$$m_\nu \sim \left(\epsilon^{|2x_l - 1|} + \epsilon^{|2x_l - 1 + \chi_{\tilde{w}}| - |\chi_{\tilde{w}}|}\right) \frac{v_u^2}{\Lambda}, \tag{2}$$

which numerically requires a large cut-off or large charges and puts obvious restrictions on the latter. Note that the wino contribution is similar or larger than the effective dimension five operator $(LH_u)^2$. Clearly, fake binos can be added to implement Type I seesaw. They contribute only to the flavor boson anomalous self-coupling.

A large lower limit on the scale Λ arises from the comparison of effective FCNC four fermion operators and the data on rare processes [11], but most results depend on the suppression of the coefficients by the flavor

[4] An obvious limitation of the FN approach with an abelian flavor symmetry is that all quantities are defined up to $O(1)$ factors. Therefore we use the symbol ”$\sim$” to express these uncertainties.

[5] Notice that the x_t parametrizes the effects of $\tan\beta$ on the determination of the χ matrices from the fermion masses. Hence only the integer part of x_t is relevant, so that $x_t = 0,\ 1$ or 2.

model. However, for broken gauged flavor symmetry there is a general limit arising from the exchange of the massive flavor gauge boson [12]. The FCNC and CPV effects appear in the flavor current when it is transformed to the physical basis because the fermion charges are different, which produces a mixing pattern comparable to the CKM matrix. This is generic and quite model independent, entailing a quite general limit on the flavor breaking scale, $\Lambda > 5 \times 10^4$ TeV. Obviously our models have to respect this bound, but then all their operators are consistent with the FCNC and CPV data without further restrictions.

Let us now turn to the discussion of the anomalies and their cancellation through the heavy states. They are extensively discussed in the literature [9] within the SM or the MSSM fermion content. Three are linear in the flavor charges, the fourth is quadratic. We display their expressions in a convenient reshuffled form as follows:

$$3\chi_{\tilde{g}} = Tr(\chi_u + \chi_d) \qquad \chi_{\tilde{w}} + \chi_{\tilde{h}} - 4\chi_{\tilde{g}} + \chi_E = -Tr(\chi_l - \chi_d)\,, \qquad 3\sum_i {x^i}_q + 3x_l + 2\chi_{\tilde{w}} + \chi_{\tilde{h}} = 1\,,$$

$$\chi_{\tilde{h}}(x_{\tilde{h}_u} - x_{\tilde{h}_d}) - \chi_E(x_E - x_{\bar{E}}) - \mathcal{F} = \sum_i [2\chi^i_u(x^i_q - x^i_u) - \chi^i_d(x^i_q - x^i_d) - \chi^i_l(x^i_l - x^i_e)]\,,$$

$$\mathcal{F} = \frac{1}{9}[Tr\{2\chi_u - \chi_d + 9\chi_l\} \pm 3\chi_{\tilde{h}}. \qquad \chi_{\tilde{g}} = 2x_{\tilde{g}}\,, \quad \chi_{\tilde{w}} = 2x_{\tilde{w}}\,, \quad \chi_{\tilde{h}} = x_{\tilde{h_u}} + x_{\tilde{h_d}}\,, \quad \chi_E = x_E + x_{\bar{E}}\,. \tag{3}$$

The contribution of the exotic heavy lepton E which is not present in the minimal model is also shown. It is easy to realize from these equations what is the minimal set of new fermions to compensate for the anomalies, and why, besides improving gauge coupling unification, it helps solving the anomaly equations. Actually, in the minimal model many solutions to these equations imply a very light particle, already excluded by experiments. Some solutions -of course, there are many - are shown in Tables 3 and 4 to illustrate various kinds of scenarios, which we now turn to discuss.

3 Ersatz gluinos, winos and Higgsinos

Since they transform in real representations of the SM gauge group, the heavy fermions are expected to get masses $O(\Lambda)$, suppressed by Froggatt-Nielsen factors. Namely, $m_f \sim \epsilon^{|\chi|_f|}\Lambda$, with $\Lambda > 5 \times 10^4$ TeV, and the exponents χ_f just defined for the exotic fermions. We look for solutions of the anomaly equations that predict states in the TeV region. With the small set of anomaly compensator fermions considered here, this is often the case once the solutions with too light states are discarded.

Besides these masses generated at the cut-off scale, Higgs couplings can give rise to mass terms analogous to the SM ones. This happens in particular, in the next-to-minimal model, for the system ($\tilde{h}, E,\ H$). Then one has the mass terms (in a sketchy notation):

$$a\epsilon^{|\chi_E|}\Lambda\bar{E}E + b\epsilon^{|\chi_{\tilde{h}}|}\Lambda\tilde{h}_u\tilde{h}_d + c\epsilon^{|x_{h_u}+x_{\bar{E}}|}\langle H_u\rangle\tilde{h}_u\bar{E} + d\epsilon^{|x_{h_d}+x_E|}\langle H_d\rangle\tilde{h}_d E, \tag{4}$$

with the explicit expressions for m_E, $m_{\tilde{h}}$, m_+ and m_-, respectively, where a, b, c, d are all $O(1)$. Because these states have not been observed, one expects $m_{E,\tilde{h}} \gtrsim v_u$.

Note that for fermions with masses beyond 1 TeV, the contribution to the oblique electroweak observables becomes negligible. However the contributions of the lighter states do need some attention. Indeed, the couplings $(E\tilde{h}H)$ can be potentially dangerous if E and $\tilde{h}$ are the lightest exotic states. In this case we compute the contribution to the T parameters using the mass insertion approximation. We find at the leading order, with

$\xi = 1 - m_E/m_{\tilde{h}}$,

$$\Delta T = \frac{1}{\sin^2\theta_W 16\pi^2}\frac{m_-^2 m_+^2}{M_W^2 m_{\tilde{h}}^2} I(\xi), \tag{5}$$

$$I(x) = \frac{x^{-4}}{3(x-2)^4}\left(4(x-2)x(x(156+x(x(x+56)-158))-60)+24(x-1)^2(20+x(x(24+x(x-12))-32))\ln(1-x)\right). \tag{6}$$

The function varies between $I(x) \sim 1-8$ with $I(0) = 6$. With conservative assumptions $(|x_{\tilde{h}_d} + x_E|, |x_{\tilde{h}_u} + x_{E'}|) \geq 1$, we find that the theory passes the EWPT with ease and with practically no constraints on the mass of the exotic fermions. As an illustration we found that for $|\Delta T| < 0.1$ and $(|x_{\tilde{h}_d} + x_E|, |x_{\tilde{h}_u} + x_{E'}|) = 1$ the bound is given by $m_{\tilde{h}} > 50\ GeV$. We have checked that all the examples in Tables 3 and 4 pass EWPT.

The contributions due to the wino and Higgsino to the T parameter in the context of the MSSM were first computed in [13]. They obtained an upper bound of $\Delta T < 0.09$. Instead, with the Z' chosen here, the two states cannot mix through the Higgs coupling and do not contribute to S and T.

As for their decays, at least the lightest ones must decay into four fermions (or be stable and become a dark matter candidate). When these particles get their masses around a few TeV, because the range of their four fermion interactions is given by Λ^{-1}, they are very long-lived at the scale of colliders, even without the suppression due to Froggatt-Nielsen factors. The decay modes will depend on the fermion structure of the interactions, the flavor distribution of the rates being more model dependent.

With the charges as defined in Tables 1 and 2, it is easy to construct the possible four fermion operators that are allowed by the discrete symmetry. Considering the Majorana nature of the fields, the only possibility for the gluino and wino is $Z' = 1/2$. For the Higgsino we also have two choices, $Z'(\tilde{h}_d) = \pm 3/18$. Below we summarize the possible decays:

$$\tilde{g} \to \bar{l}q\bar{u}_R,\ \ \bar{l}d_R\bar{q},\ \ \bar{e_R}\bar{u_R}d_R(+\text{h.c.}) \qquad\qquad \tilde{w} \to \bar{l}q\bar{u}_R,\ \ \bar{l}d_R\bar{q},\ \ \bar{l}\bar{l}d_R,\ \ \bar{l}\bar{H}_u(+\text{h.c.})$$

$$Z'(\tilde{h}_d) = -3/18:\ \tilde{h}_d \to \bar{q}\bar{q}\bar{q},\ \ \bar{q}\bar{u}\bar{d} \qquad Z'(\tilde{h}_d) = 3/18:\ \tilde{h}_d \to qdd \qquad\qquad E \to \bar{q}\bar{q}\bar{u},\ \ \bar{d}\bar{u}\bar{u}\,.$$

The family patterns of the decay products are very model dependent. The lifetimes determined by these decay modes of the lightest exotic particles, in some specific examples are given in the tables.

The addition of adjoint and weak doublet fields has the potential to drive gauge coupling unification. This can be anticipated from the close resemblance of these models with the split SUSY scenario. Even if exact unification is not achieved, in the majority of these models a significant improvement in gauge coupling unification should be possible. In order to make a quantitative study we define the following parameters,

$$\alpha_i(M_{\text{GUT}}^{-1} - \alpha_i(M_Z)^{-1} = \Delta_i = \Delta_i^{\text{SM}} + \Delta_i^{\text{new}}, \tag{7}$$

where i corresponds to the three gauge groups of the SM and Δ_i^{SM} and Δ_i^{new} are the contributions from the SM (new) states , respectively, to the running of the gauge couplings from M_Z to M_{GUT}. Using the experimentally measured low energy values of the gauge couplings, the condition for gauge coupling unification in a generic model reduces to the following expressions [14],

$$\rho = \frac{\Delta_2 - \Delta_3}{\Delta_1 - \Delta_2} = 0.719 \pm 0.005\,, \qquad\qquad \Delta_1 - \Delta_2 = 29.42 \pm 0.03. \tag{8}$$

To analyze the improvement in gauge coupling unification quantitatively we device the following strategy: for every model a M_{GUT} value is defined from the $\Delta_1 - \Delta_2$ given in Eqs. 8. Then the following parameter is computed for the new model,

$$r = \frac{|\rho_{\text{new}} - 0.719|}{|\rho_{\text{SM}} - 0.719|}. \tag{9}$$

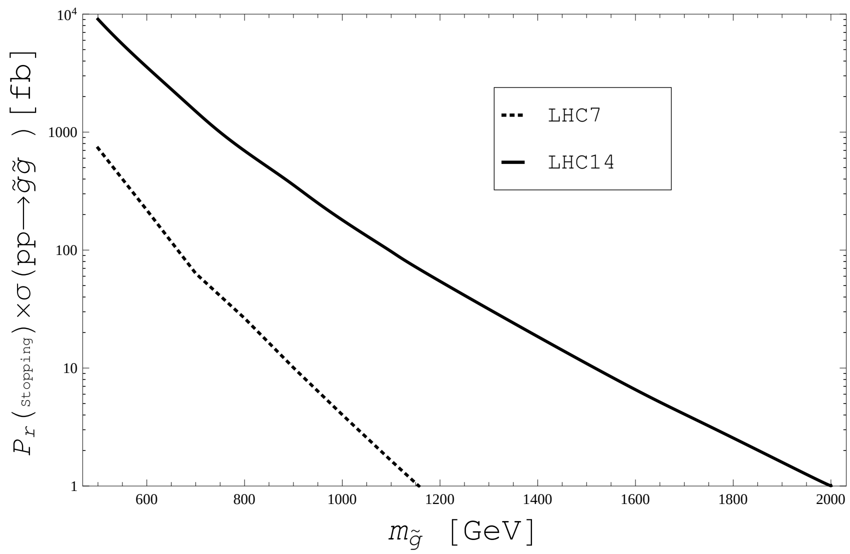

Figure 1: *The stopping cross-section for the gluinos at LHC.*

In Tables 3 and 4 this parameter is quoted for the specific models and it is also indicated when the unification is consistent with the present experimental errors.

4 Experimental consequences

The exotic fermions introduced here, have four fermion interactions that are determined by the ad hoc but motivated assignments of the Z' charges as given in Table 2. Deviation from these specific choices can summarily kill all possible interaction leaving these fermions stable at the cosmological scale and making them possible dark matter candidates. For the Majorana fermions the choices for the Z' assignment are restricted to $\pm 1/2$ or 0. An interesting scenario arises when $\chi_{\tilde{\omega}}$ is even. In this case the Z' charge for the exotic wino is zero and therefore it cannot couple to the SM fermions. If this is the lightest exotic fermion with mass around $\sim 2.7\,TeV$ it can be a viable Minimal Dark Matter (MDM) candidate [15] that satisfies the 7 year WMAP constraints [16]. However the nuclear cross-sections of these species get an unavoidable Sommerfeld enhancement putting them in imminent danger of being ruled out by indirect searches, specially by the anti-proton data [17]. At present the experimental observations are subject to large astrophysical uncertainties leaving enough room for the wino to survive as a DM candidate.

Another interesting scenario would be to consider the Z' charge of exotic Higgsino chosen to make it stable. It is known that a pure Higgsino like stable particle can satisfy the WMAP data if it has a mass around $\sim 1\,TeV$ [18]. Unfortunately the coupling of the Higgsino to the Z boson makes it unfavorable from indirect searches.

The exotic fermions introduced can be within the experimental reach of the LHC as is evident from Tables 3 and 4. These exotic particles can be produced at the collider and all of them are metastable except the exotic wino which decays easily via the interaction $\tilde{w}H_uL$. The signals of the SU(2) adjoint triplet at the colliders closely resemble the scenario of generic Type III seesaw models. The triplet is produced by gauge interactions $q\bar{q} \to \tilde{w}^0\tilde{w}^0, \tilde{w}^+\tilde{w}^-$ and $u\bar{d} \to \tilde{w}^+\tilde{w}^0$. The latter is expected to have the highest cross-section at the LHC

which is $\sim 35fb$ for $m_{\tilde{w}} = 500GeV$ and it falls to $\sim 1fb$ when the mass reaches $1TeV$. Once produced they mainly decay by $\tilde{w}^0 \to \nu h$ and $\tilde{w}^{\pm} \to l^{\pm}h$. For the range of parameters where Eq. 2 gives the correct neutrino mass one expects a displaced vertex. See [19] for a detailed study of signals in Type III seesaw models at LHC, that closely resemble the scenario where the exotic wino is the only observable particle in the model.

Consider a scenario where the long living exotic gluino is the lightest particle within the reach of the LHC. In this case the exotic gluinos can be copiously produced at the hadron collider. Once produced these particles will hadronize into R-hadrons and R-mesons. They will lose considerable amount of energy as they travel through the detector. A fraction of these hadronized exotic gluinos will come to rest within the detector and then decay at a later instance [20] giving rise to an interesting signature. Both the CMS [6] and ATLAS [7] have already published their results on long living gluinos. The present bound on the particle masses is around $400\ GeV$, however we note that most of these studies were carried out within the framework of the split SUSY scenario. In [6] the probability for a produced gluino to be stopped within the CMS detector was simulated. In Figure 1 we show the cross-section of stopped gluinos as a function of gluino mass. We have used CalcHEP 2.3.5 [21] to calculate the cross-section. Once stopped these metastable particles will decay generally in the quiescent period, i.e., out of sync with the proton-proton collision at the collider. The decay will lead to a signal like: $1\ prompt\ lepton + 2\ jets$. This decay is different from the ones studied at the LHC experiments within the split SUSY paradigm. We note that the prompt lepton in the decay of the exotic gluino in this class of models clearly distinguishes it from the split SUSY signals.

The production cross-section for the exotic Higgsino and the exotic lepton is much smaller than that of the exotic gluino. Nevertheless their decay after being captured in the detector can give rise to spectacular signals that can easily be identified. However it is their slow production rates that makes it difficult to observe at low energy/luminosity.

5 Conclusion

In this paper, we have studied the properties of particles with the quantum numbers of the sfermions of supersymmetric versions of the SM in the total absence of supersymmetry. We have shown that they can play a role in gauged flavor symmetry to cancel the anomalies due to the quarks and leptons. A crucial point is the discrete symmetry that survives flavor symmetry breaking - proton hexality - to stabilize the flavor models (and the proton!). In most of the viable examples of this framework, some of the new states might appear within the reach of the LHC as metastable particles with characteristic decay patterns. Alternatively, the set up encompasses a version of the Type III seesaw phenomenology, with a weak isospin triplet fermion whose mass is light because of the proposed mechanism. Another consistent possibility is that this triplet is stable and is a dark matter candidate. This scenario can be easily extended to non-abelian flavor symmetry.

Acknowledgments: We thank Marco Cirelli for discussions. This work has been partially supported by the Agence Nationale de la Recherche under contract ANR 2010 BLANC 041301. The work of TSR is supported by EU ITN , contract ”UNILHC” PITN-GA-2009-237920, the CEA-Eurotalents program. HS is supported by a CAPES Foundation (Ministry of Education of Brazil) Postdoc Fellowship.

Minimal Model					
Model Parameters		$Exotic\ fields$	$Lightest\ particle$	r	$Decay\ width$
$Sl.No.$	$SM\ fields$				
1	$x_q = (2, 0, -2)$ $x_{\bar{u}} = (6, 4, 2)$ $x_{\bar{d}} = (7, 7, 7)$ $x_t = 0,\ x_H = -2$ $x_l = 4,\ z'(\tilde{h}_d) = -\frac{3}{18}$ $x_{\bar{e}} = (-9, 3, -1)$	$\chi_{\tilde{g}} = -11$ $\chi_{\tilde{w}} = -1$ $\chi_{\tilde{h}} = -7$	$m_{\tilde{g}} \sim 1.5\ TeV$	0.33	$C\tau(\tilde{g}_1) \sim 5.6 \times 10^{12}\ km$
2	$x_q = (7, 5, 3)$ $x_{\bar{u}} = (1, -1, -3)$ $x_{\bar{d}} = (6, 6, 6)$ $x_t = 1,\ x_H = 5$ $x_l = -3,\ z'(\tilde{h}_d) = -\frac{3}{18}$ $x_{\bar{e}} = (4, 2, -4)$	$\chi_{\tilde{g}} = -3$ $\chi_{\tilde{w}} = -11$ $\chi_{\tilde{h}} = 9$	$m_{\tilde{w}} \sim 0.5\ TeV$	1.48	$\tilde{w}$ is unstable
3	$x_q = (2, 0, -2)$ $x_{\bar{u}} = (6, 4, 2)$ $x_{\bar{d}} = (-1, -1, -1)$ $x_t = 0,\ x_H = 6$ $x_l = -4,\ z'(\tilde{h}_d) = -\frac{3}{18}$ $x_{\bar{e}} = (5, -7, -5)$	$\chi_{\tilde{g}} = -3$ $\chi_{\tilde{w}} = 1$ $\chi_{\tilde{h}} = 9$	$m_{\tilde{h}} \sim 0.5\ TeV$	0.33	$C\tau(\tilde{h}_d) \sim 8.2 \times 10^{6}\ km$ $C\tau(\tilde{h}_u) \sim 4.4 \times 10^{8}\ km$
4	$x_q = (5, 3, 1)$ $x_{\bar{u}} = (3, 1, -1)$ $x_{\bar{d}} = (-2, -2, -2)$ $x_t = 2,\ x_H = 2$ $x_l = 4,\ z'(\tilde{h}_d) = -\frac{3}{18}$ $x_{\bar{e}} = (-1, -3, -7)$	$\chi_{\tilde{g}} = -5$ $\chi_{\tilde{w}} = -10$ $\chi_{\tilde{h}} = 2$	$m_{\tilde{w}} \sim 1.5\ TeV$	0.86	$\tilde{w}$ is a dark matter candidate

Table 3: Examples for the minimal model

Next to Minimal Model					
Model Parameter		$Exotic\ fields$	$Lightest\ particle$	r	$Decay\ width$
$Sl.No.$	$SM\ fields$				
1	$x_q = (4,2,0)$ $x_{\bar{u}} = (4,2,0)$ $x_{\bar{d}} = (-1,-1,-1)$ $x_t = 0,\ x_H = 4$ $x_l = -4,\ z'(\tilde{h}_d) = -\frac{3}{18}$ $x_{\bar{e}} = (-7,-5,3)$	$\chi_{\tilde{g}} = -5$ $\chi_{\tilde{w}} = -7$ $\chi_{\tilde{h}} = 9$ $\chi_E = 2$	$m_{\tilde{h}} \sim 1.5\,TeV$	0.09 Unification	$C\tau(\tilde{h}_d) \sim 7.5 \times 10^7\ km$ $C\tau(\tilde{h}_u) \sim 1.8 \times 10^9\ km$
2	$x_q = (6,4,2)$ $x_{\bar{u}} = (2,0,-2)$ $x_{\bar{d}} = (-3,-3,-3)$ $x_t = 2,\ x_H = 2$ $x_l = -4,\ z'(\tilde{h}_d) = -\frac{3}{18}$ $x_{\bar{e}} = (-3,5,3)$	$\chi_{\tilde{g}} = -5$ $\chi_{\tilde{w}} = -9$ $\chi_{\tilde{h}} = -5$ $\chi_E = 4$	$m_{\tilde{w}} \sim 1.5\,TeV$	0.09 Unification	$\tilde{w}$ in unstable
3	$x_q = (4,2,0)$ $x_{\bar{u}} = (4,2,0)$ $x_{\bar{d}} = (1,1,1)$ $x_t = 2,\ x_H = 0$ $x_l = 4,\ z'(\tilde{h}_d) = -\frac{3}{18}$ $x_{\bar{e}} = (-9,-1,-3)$	$\chi_{\tilde{g}} = -7$ $\chi_{\tilde{w}} = -10$ $\chi_{\tilde{h}} = -9$ $\chi_E = 1$	$m_{\tilde{w}} \sim 1.6\,TeV$	0.05 Unification	$\tilde{w}$ is a dark matter candidate
4	$x_q = (-1,-3,-5)$ $x_{\bar{u}} = (9,7,5)$ $x_{\bar{d}} = (0,0,0)$ $x_t = 0,\ x_H = 8$ $x_l = 6,\ z'(\tilde{h}_d) = -\frac{3}{18}$ $x_{\bar{e}} = (-7,-9,-11)$	$\chi_{\tilde{g}} = -1$ $\chi_{\tilde{w}} = 7$ $\chi_{\tilde{h}} = -4$ $\chi_E = -7$	$m_E \sim m_{\tilde{w}} \sim 0.5\,TeV$	0.07 Unification	$\tilde{w}$ is unstable $C\tau(E) \sim 959\ km$ $C\tau(\bar{E}) \sim 1.1 \times 10^9\ km$
5	$x_q = (5,3,1)$ $x_{\bar{u}} = (3,1,-1)$ $x_{\bar{d}} = (2,2,2)$ $x_t = 0,\ x_H = 0$ $x_l = -4,\ z'(\tilde{h}_d) = -\frac{3}{18}$ $x_{\bar{e}} = (-3,-1,1)$	$\chi_{\tilde{g}} = -9$ $\chi_{\tilde{w}} = -3$ $\chi_{\tilde{h}} = -8$ $\chi_E = 5$	$m_{\tilde{g}} \sim 1.5\,TeV$	0.9	$C\tau(\tilde{g}) \sim 7.1 \times 10^6\ km$

Table 4: Examples for the next to minimal model

References

[1] C. D. Froggatt and H. B. Nielsen, Nucl. Phys. B **147** (1979) 277.

[2] C. A. Savoy, M. Thormeier, JHEP **1103** (2011) 128. [arXiv:1003.2090 [hep-ph]].

[3] H. K. Dreiner, C. Luhn, M. Thormeier, Phys. Rev. **D73** (2006) 075007. [hep-ph/0512163].

[4] N. Arkani-Hamed and S. Dimopoulos, JHEP **0506** (2005) 073 [hep-th/0405159].

[5] N. Arkani-Hamed, S. Dimopoulos, G. F. Giudice and A. Romanino, Nucl. Phys. B **709** (2005) 3 [hep-ph/0409232].

[6] V. Khachatryan *et al.* [CMS Collaboration], Phys. Rev. Lett. **106** (2011) 011801 [arXiv:1011.5861 [hep-ex]]; V. Khachatryan *et al.* [CMS Collaboration], JHEP **1103** (2011) 024 [arXiv:1101.1645 [hep-ex]].

[7] G. Aad *et al.* [ATLAS Collaboration], Phys. Lett. B **703** (2011) 428 [arXiv:1106.4495 [hep-ex]]; G. Aad *et al.* [ATLAS Collaboration], Phys. Lett. B **701** (2011) 1 [arXiv:1103.1984 [hep-ex]]; G. Aad *et al.* [ATLAS Collaboration], arXiv:1201.5595 [hep-ex].

[8] O. J. P. Eboli, C. A. Savoy and R. Z. Funchal, arXiv:1112.5108 [hep-ph].

[9] L. E. Ibanez and G. G. Ross, Phys. Lett. B **332** (1994) 100 [hep-ph/9403338]; P. Binetruy and P. Ramond, Phys. Lett. B **350** (1995) 49 [hep-ph/9412385]; E. Dudas, S. Pokorski and C. A. Savoy, Phys. Lett. B **356** (1995) 45 [hep-ph/9504292]; E. Dudas, S. Pokorski and C. A. Savoy, Phys. Lett. B **369** (1996) 255 [hep-ph/9509410]; P. Binetruy, S. Lavignac and P. Ramond, Nucl. Phys. B **477** (1996) 353 [hep-ph/9601243].

[10] R. Foot, H. Lew, X. G. He and G. C. Joshi, Z. Phys. C **44** (1989) 441.

[11] G. Isidori, Y. Nir and G. Perez, Ann. Rev. Nucl. Part. Sci. **60** (2010) 355 [arXiv:1002.0900 [hep-ph]].

[12] N. Arkani-Hamed, L. J. Hall, D. Tucker-Smith and N. Weiner, Phys. Rev. D **61** (2000) 116003 [hep-ph/9909326].

[13] R. Barbieri, L. Maiani, Nucl. Phys. **B224** (1983) 32.

[14] P. Nath and P. Fileviez Perez, Phys. Rept. **441** (2007) 191 [hep-ph/0601023].

[15] M. Cirelli, N. Fornengo, A. Strumia, Nucl. Phys. **B753** (2006) 178-194. [hep-ph/0512090].

[16] E. Komatsu *et al.* [WMAP Collaboration], Astrophys. J. Suppl. **192** (2011) 18 [arXiv:1001.4538 [astro-ph.CO]].

[17] M. Cirelli, R. Franceschini and A. Strumia, Nucl. Phys. B **800** (2008) 204 [arXiv:0802.3378 [hep-ph]]; M. Cirelli, G. Corcella, A. Hektor, G. Hutsi, M. Kadastik, P. Panci, M. Raidal and F. Sala *et al.*, JCAP **1103** (2011) 051 [arXiv:1012.4515 [hep-ph]].

[18] G. F. Giudice and A. Romanino, Nucl. Phys. B **699** (2004) 65 [Erratum-ibid. B **706** (2005) 65] [arXiv:hep-ph/0406088].

[19] R. Franceschini, T. Hambye, A. Strumia, Phys. Rev. **D78** (2008) 033002. [arXiv:0805.1613 [hep-ph]].

[20] A. Arvanitaki, S. Dimopoulos, A. Pierce, S. Rajendran, J. G. Wacker, Phys. Rev. **D76** (2007) 055007. [hep-ph/0506242].

[21] A. Pukhov, hep-ph/0412191.